$SO(10)$ GUTs with Gauge Mediated Supersymmetry Breaking

M. Frank[a], H. Hamidian[b], and K. Puolamäki[c]
[a]Department of Physics, Concordia University, 1455 De Maisonneuve Blvd. W., Montréal, PQ H3G 1M8, Canada
[b]Department of Physics, University of Illinois at Chicago, 845 W. Taylor Street, Chicago, IL 60607, USA
[c]Helsinki Institute of Physics, P.O.Box 9, FIN-00014 University of Helsinki, Finland

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We explore the phenomenology of supersymmetric $SO(10)$ grand unified theories with gauge mediated supersymmetry breaking. We show that if $SO(10)$ breaking proceeds through intermediate left-right symmetric gauge groups which are broken at the supersymmetry breaking scale, then perturbative unification allows the existence of only a few consistent models with very similar phenomenological consequences. We list and discuss some distinctive signatures of these theories. The most remarkable feature of the class of theories introduced here is that, unlike in models with simpler symmetry breaking chains, the set of allowed messengers is practically unique.

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Introduction. If the standard model (SM) indeed descends from a supersymmetric theory, a number of criteria must be met and some major questions need to be answered. The most important question among these is to understand the origin and nature of supersymmetry (SUSY) breaking. The currently popular models include gravity-mediated SUSY breaking, gauge-mediated SUSY breaking, and $U(1)$-mediated SUSY breaking [1]. Each of these models exhibits certain desirable features and drawbacks of its own; however, the main underlying assumption, i.e. the existence of a hidden sector where SUSY is broken and the means by which the breaking is transmitted to the visible (low-energy) world of the SM, is shared by all of these approaches. In theories with gauge mediated SUSY breaking (GMSB) supersymmetry-breaking soft terms are generated by a set of particles, called messengers, at a scale Λ_M (called the messenger scale) which is *a priori* unrelated to the GUT scale [2]. A distinguishing and attractive feature of GMSB models is that they naturally lead to degenerate squark and slepton masses and thus alleviate the flavor changing neutral current (FCNC) problem of the minimal supersymmetric standard model (MSSM). Furthermore, GMSB theories are highly predictive; they lead to a dramatic reduction of the number of free parameters and their predictions will be testable in the not-too-distant future [3].

In GMSB theories SUSY breaking in the (unspecified) hidden sector is communicated to the visible sector through the $SU(3)_C \times SU(2)_L \times U(1)_Y$ SM gauge interactions of the "messenger" fields with the visible sector. In the minimal version of GMSB, the messenger fields belong to the $\mathbf{5} + \overline{\mathbf{5}}$ or $\mathbf{10} + \overline{\mathbf{10}}$ representations of the $SU(5)$ gauge group and there exists at least one singlet superfield S which couples to vector-like messenger superfields $V + \overline{V}$ through the superpotential interaction

$$W_{mess} = \lambda_V S V \overline{V}, \tag{1}$$

where the Yukawa couplings, λ_V's in (1), are assumed to coincide at the unification scale M_G. Consequently, the spectrum at the messenger scale consists of a set of fields in complete $SU(5)$ representations and the mass splitting among the fields can be determined from the renormalization group (RG) running of the messenger Yukawa couplings from M_G down to the messenger scale Λ_M. The (generalized) non-minimal versions of GMSB theories, in which the messenger fields do not necessarily form complete $SU(5)$ GUT multiplets, have also been studied by a number of authors [4–6]. These studies indicate that GMSB theories based on $SU(5)$ GUTs are phenomenologically disfavored. However, since the high predictability and simplicity of GMSB theories are hard to achieve otherwise, it is important to study other (SUSY) GUT gauge groups which may lead to more realistic models.

An attempt to go beyond the minimal GMSB gauge group was made in [7] where the authors embedded the electroweak gauge group in the gauge group $SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$ [or $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$] at the SUSY breaking scale Λ_{SUSY}. The distinguishing features of these models include the automatic conservation of R-parity, non-vanishing neutrino masses, and a unified hidden-plus-messenger sector potential. This has led us to investigate whether $SO(10)$ SUSY GUTs (which contain these gauge groups) with GMSB could lead to phenomenologically attractive scenarios. It is the purpose of this Letter to show that not only does this seem to be the case, but also that there are only a few scenarios (with very similar testable predictions) which are consistent with all the constraints.

$SO(10)$ *Models and GMSB.* The usual model building in GMSB theories based on the $SU(5)$ gauge group suffers from a number of serious drawbacks and leaves some important questions unanswered. In particular, one encounters problems related to the nucleon decay rates, lack of a natural mechanism to generate neutrino masses, the existence of arbitrary R-parity violating interactions, and the SUSY CP problem. A viable alternative for naturally avoiding most of these problems is to consider $SO(10)$ grand unification instead. In $SO(10)$ GUTs the dangerous colour triplet

Higgs boson can naturally be made very heavy—thus suppressing rapid proton decay —and by incorporating left-right (LR) supersymmetric theories, which can arise through $SO(10)$ breaking, the SUSY CP problem can be solved while automatically conserving R-parity [8]. Furthermore, as an additional bonus, since a full generation of fermions comes in one spinorial representation which includes right-handed neutrinos, the see-saw mechanism can be used to naturally generate small neutrino masses [9].

$SO(10)$ symmetry breaking can proceed in essentially two ways: $SO(10)$ can break down to $SU(5) \times U(1)$ at $\sim 10^{18}$ GeV with a further breaking down to the MSSM at $\sim 10^{16}$ GeV. Alternatively, $SO(10)$ can break down to some LR symmetric gauge group G_{LR} (such as $G_{224} = SU(2)_L \times SU(2)_R \times SO(4)_C$) with a subsequent breaking down to the MSSM at some intermediate scale. Here we assume that $SO(10)$ first breaks down to an intermediate LR symmetric group G_{LR} at the scale M_G, followed by the breaking of G_{LR} at the scale M_R down to the MSSM. The scale M_G lies below the Planck scale $\sim 10^{19}$ GeV and must be no less than 10^{16} GeV to ensure nucleon stability. To keep things as simple as possible, we further assume that there are no other symmetry breaking scales between M_R and M_G.

The supersymmetric ($SO(10)$-based) LR models described here have the gauge groups: $G^I_{LR} = SU(2)_L \times U(1)_{I_{3R}} \times SU(3)_C \times U(1)_{B-L}$ and $G^{II}_{LR} = SU(2)_L \times SU(2)_R \times SU(3)_C \times U(1)_{B-L}$, and further assumptions are needed to include SUSY breaking in these models. Since chirality plays a very important role in SUSY it is very natural to expect SUSY and LR symmetry breaking scales to be somehow related. A very attractive possibility is to simply take $\Lambda_{SUSY} = M_R$ which not only connects the SUSY breaking and the gauge symmetry breaking scales, but also requires that the electroweak symmetry breaking remain radiative. A consequence of this assumption is that $\Lambda_{SUSY} \sim M_R \sim 100$ TeV, which follows from the usual (MSSM) requirement that the sparticle masses stay in the TeV range. We will thus limit the ranges of the left-right breaking and the GUT scales as:

$$10^5 \text{ GeV} < M_R < 10^7 \text{ GeV and } 10^{16} \text{ GeV} < M_G < 10^{19} \text{ GeV}. \tag{2}$$

Furthermore, as is usually done in the study of SUSY GUTs, we assume that the gauge couplings unify at M_G and remain perturbative up to the Planck scale. Putting all these ingredients together we now proceed by specifying the LR models studied here.

Model I: $SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$. This model is phenomenologically interesting for a number of reasons: It contains all the usual matter multiplets *plus* right-handed neutrinos and (due to the absence of baryon-number violating terms in the superpotential) automatically conserves R-parity. The superpotential which describes this model is:

$$\begin{aligned} W &= h_u Q H_u u^c + h_d Q H_d d^c + h_e L H_d e^c + h_\nu L H_u \nu^c \\ &+ \mu H_u H_d + f \delta \nu^c \nu^c + M_R \delta \bar{\delta} + W_m, \end{aligned} \tag{3}$$

where W_m is the messenger sector superpotential. The particle content of this model consists of: The doublets $Q(2,0,1/6)$, $L(2,0,-1/2)$, the singlets $u^c(1,-1/2,-1/6)$, $d^c(1,1/2,-1/6)$, $e^c(1,1/2,1/2)$, $\nu^c(1,-1/2,1/2)$, the Higgs doublets $H_u(2,1/2,0)$ and $H_d(2,-1/2,0)$, which are the same as in the MSSM, two Higgs triplets, $\delta(1,1,-1)$ and $\bar{\delta}(1,-1,1)$, which break the $U(1)_{I_{3R}} \times U(1)_{B-L}$ symmetry down to $U(1)_Y$ of the standard model, the gauge bosons $W(3,0,0)$, $B(1,0,0)$, and $V(1,0,0)$, and the superpartners of all these fields.

Model II: $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. The phenomenology of the LR supersymmetric model based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ has been extensively studied in recent years [10]. Among other attractive features, this model simultaneously solves both the strong and weak CP problems and can also accomodate automatic conservation of R-parity. This model can be specified by the superpotential:

$$\begin{aligned} W &= \mathbf{h}_q^{(i)} Q_L^T \tau_2 \Phi_i \tau_2 Q_R + \mathbf{h}_l^{(i)} L_L^T \tau_2 \Phi_i \tau_2 L_R \\ &+ i(\mathbf{h}_{LR} L_L^T \tau_2 \delta_L L_L + \mathbf{h}_{LR} L_R^T \tau_2 \Delta_R L_R) \\ &+ M_{LR} \left[\mathrm{Tr}(\Delta_L \delta_L + \mathrm{Tr}(\Delta_R \delta_R) \right] \\ &+ \mu_{ij} \mathrm{Tr}(\tau_2 \Phi_i^T \tau_2 \Phi_j) + W_m, \end{aligned} \tag{4}$$

where W_m denotes the messenger sector superpotential. The particle content of this model consists of the doublets $Q_L(2,1,1/6)$, $L_L(2,1,-1/2)$, $Q_R(1,2,-1/6)$, $L_R(1,2,1/2)$, the bi-doublet Higgs fields $\Phi_u(2,2,0)$ and $\Phi_d(2,2,0)$, the triplet Higgs fields $\Delta_L(3,1,-1)$, $\Delta_R(1,3,-1)$, $\delta_L(3,1,1)$, and $\delta_R(1,3,1)$, the gauge bosons $W_L(3,1,0)$, $W_R(1,3,0)$, and $V(1,1,0)$, and the corresponding superpartners.

In both models I and II the relation between the $U(1)_{B-L}$ gauge coupling α_{B-L} and the GUT-normalized gauge coupling α_V is fixed through $\alpha_V = \frac{2}{3}\alpha_{B-L}$. The messenger sector in both models is described by N_f flavours of chiral superfields Φ_i and $\overline{\Phi}_i$ ($i = 1, \cdots, N-f$) which belong to the $\mathbf{r} + \overline{\mathbf{r}}$ representation of the corresponding gauge group.

If one assumes gauge coupling constant unification and requires that the messengers form complete GUT multiplets, then the presence of intermediate scale messenger fields leaves M_G unchanged. In this case, due to the contribution of the messenger fields, one obtains [2]:

$$\delta\alpha_{GUT}^{-1} = -\frac{N}{2\pi}\ln\frac{M_G}{\Lambda_M}, \tag{5}$$

where $N = \sum_{i=1}^{N_f} n_i$, with n_i denoting twice the Dynkin index of the $\mathbf{r}$ representation for the ith flavor. This leads to the constraint:

$$N \lesssim 150/\ln\frac{M_G}{\Lambda_M}, \tag{6}$$

as a result of the perturbativity hypothesis at the GUT scale.

To search for possible GMSB scenarios that are consistent with the assumptions and constraints that we have invoked so far, let us begin by listing the messenger fields. In choosing messenger fields, which form the messenger sector, we maintain the constraint that they should occupy the same representation as Model I or Model II chiral superfields. (The motivation is that stable particles with exotic quantum numbers are a disaster for cosmology [4]). The possible messenger fields in Model I, which transform under the gauge group $SU(3)_C \times SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$, are given by: $Q_8 = (8,1,0,0), L_3 = (1,3,0,0), \Delta + \overline{\Delta} = (1,3,0,-1) + conj., \Delta^c + \overline{\Delta^c} = (1,1,-1,1) + conj., H + \overline{H} = (1,2,\frac{1}{2},0) + conj., Q + \overline{Q} = (3,2,0,\frac{1}{6}) + conj., U^c + \overline{U^c} = (\overline{3},1,-\frac{1}{2},-\frac{1}{6}) + conj., D^c + \overline{D^c} = (\overline{3},1,\frac{1}{2},-\frac{1}{6}) + conj., L + \overline{L} = (1,2,0,-\frac{1}{2}) + conj., e^c + \overline{e^c} = (1,1,\frac{1}{2},\frac{1}{2}) + conj., \nu^c + \overline{\nu^c} = (1,1,-\frac{1}{2},\frac{1}{2}) + conj.$, and those of Model II, transforming under the $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge group, are given by: $Q_8 = (8,1,1,0), Q_3 = (1,3,1,0), Q_3^c = (1,1,3,0), \phi = (1,2,2,0), Q + \overline{Q} = (3,2,1,\frac{1}{6}) + conj., Q^c + \overline{Q^c} = (3,1,2,-\frac{1}{6}) + conj., L + \overline{L} = (1,2,1,-\frac{1}{2}) + conj., L^c + \overline{L^c} = (1,1,2,\frac{1}{2}) + conj., \Delta + \overline{\Delta} = (1,3,1,-1) + conj., \Delta^c + \overline{\Delta^c} = (1,1,3,1) + conj.$.

To restrict the messenger sector by using the unification requirement the RG β-functions and the matching conditions must be used. At the LR breaking scale M_R the couplings are required to match those of the MSSM. Denoting the β-functions of the $SU(2)_L$, $SU(2)_R$, $U(1)_{B-L}$, and $SU(3)_C$ gauge groups respectively by β_L, β_R, β_V, and β_C, the one-loop RG equations at the scale M_R are:

$$\alpha_I^{-1}(M_R) = \alpha_G^{-1} + \beta_I(t_G - t_R), \tag{7}$$

where $t_R = \frac{1}{2\pi}\ln\frac{M_R}{M_Z}$, $t_G = \frac{1}{2\pi}\ln\frac{M_G}{M_Z}$, and $I = L, R, V, C$.

The one-loop matching conditions at the scale M_R read as follows:

$$\begin{aligned} \frac{5}{3}\alpha_1^{-1}(M_R) &= \alpha_R^{-1}(M_R) + \frac{2}{3}\alpha_V^{-1}(M_R), \\ \alpha_2^{-1}(M_R) &= \alpha_L^{-1}(M_R), \\ \alpha_3^{-1}(M_R) &= \alpha_C^{-1}(M_R), \end{aligned} \tag{8}$$

where α_1, α_2, and α_3 correspond to the $U(1)_Y$, $SU(2)_L$, and $SU(3)_C$ gauge couplings respectively. Combining equations (7) and (8) then yields:

$$\alpha_k^{-1}(M_Z) = \alpha_G^{-1} + \beta_k^{MSSM}(t_R) + \beta_k^{LR}(t_G - t_R), \tag{9}$$

with $k = 1, 2, 3$ and

$$\beta^{LR} = \begin{pmatrix} \beta_1^{LR} \\ \beta_2^{LR} \\ \beta_3^{LR} \end{pmatrix} = \begin{pmatrix} \frac{3}{5}\beta_R + \frac{2}{5}\beta_V \\ \beta_L \\ \beta_C \end{pmatrix}, \tag{10}$$

where the MSSM β-functions $\beta_k^{MSSM} = (33/5, 1, -3)^T$ have been used.

By eliminating α_G from equation (9) it is straightforward to obtain the following limits on the differences between the LR β functions:

$$3.1 < \beta_2^{LR} - \beta_3^{LR} < 4.1 \text{ and } 7.4 < \beta_1^{LR} - \beta_3^{LR} < 9.9, \tag{11}$$

and demanding that the gauge couplings remain perturbative all the way up to the Planck scale equation (9) yields:

$$\beta_1^{LR} < 10.4\ ,\ \beta_2^{LR} < 6.1 \text{ and } \beta_3^{LR} < 3.0, \tag{12}$$

which constrains the number of messenger fields.

By examining the β functions for Models I and II along with the constraints (11) and (12) we find that there are no consistent solutions in Model II. Model I, on the other hand, leads to consistent solutions with the messenger multiplicities:

$$n_8 = n_3 = n_H + n_L = 1 \text{ and } n_{e^c} + n_{\nu^c} = 0, 1. \tag{13}$$

According to (13) the messenger sector in Model I consists of a color octet ($n_8 = 1$), an $SU(2)_L$ triplet ($n_3 = 1$), a pair of H or L type messenger fields ($n_H + n_L = 1$), and a pair of e^c or ν^c type fields ($n_{e^c} + n_{\nu^c} = 0, 1$). There are thus a total of six solutions for the messenger multiplicities, all of which (as will be described in more detail elsewhere [11]) lead to very similar mass parameters for the MSSM and render our scheme extremely predictive. In all these cases the $SU(2)_L$ and $SU(3)_C$ gauge couplings meet at $M_G \simeq 2.0 \times 10^{16}$ GeV for 10^5 GeV $< M_R < 10^7$ GeV. For the solution with $n_{e^c} + n_{\nu^c} = 0$, α_1 and α_3 meet at $M'_G \simeq 1.9 \times 10^{16}$ GeV, which is within 6% of the GUT scale. When $n_{e^c} + n_{\nu^c} = 1$ the mismatch is much worse; in this case α_1 and α_3 meet at $M'_G \simeq 4 - 5 \times 10^{15}$ GeV. However, since the mismatch can be attributed to threshold effects, we will consider this solution as well. Another constraint on the models studied here comes from the ratio $\tan^2\theta_R = \alpha_{B-L}(M_R)/\alpha_R(M_R)$, which is a very important parameter in LR theories. It is completely determined by the unification condition and the chosen messenger multiplicities; in our case one has $1.3 \leq \tan^2\theta_R \leq 1.6$, as shown in table I. To compute the sparticle mass spectrum for the models studied here this value must be taken into account.

The sparticle spectrum. The messenger masses in the LR models described here can be calculated once the messenger sector and the messenger scale are fixed. By matching the LR model and the MSSM the values of the MSSM parameters at the messenger scale and the RG β-functions can be used to calculate the full MSSM particle spectrum as a function of $\tan\beta$ [5]. One could further reduce the number of parameters and solve the supersymmetric CP-problem, as was done in [6], by requiring that the bilinear scalar coupling B vanish at the messenger scale; a condition which fixes $\tan\beta$. Here, however, we wish to consider a broader possibility by letting $\tan\beta$ (or equivalently the bilinear scalar coupling) remain a free parameter. The resulting sparticle spectrum—calculated in terms of $\Lambda_{SUSY}, \Lambda_M$, and $\tan\beta$—can generally be divided into a light and a heavy sector. In our case the light sector consists of the lightest neutralino ($\tilde{\chi}^0_1$), the chargino ($\tilde{\chi}^\pm_1$), and the light slepton mass eigenstates ($\tilde{e}_1$, $\tilde{\mu}_1$, $\tilde{\tau}_1$), while the squarks comprise the heavy sector.

An important issue which requires particular care when calculating the sparticle spectrum is radiative gauge symmetry breaking. To ensure radiative (gauge) symmetry breaking, one must check that the resulting vacuum is physical. This can be achieved if all the mass-squared eigenvalues of the charged scalars remain positive and above the current experimental limits. We have taken these important constraints into account in our calculations of the sparticle spectrum (a complete listing of which will be given in [11]). Since the heavy top squark drives the radiative symmetry breaking, we find that for all values of parameters the mass-squared term of the up-type Higgs boson acquires a negative value and thus always leads to radiative symmetry breaking. [As is usual in GMSB models, the lightest sparticle turns out to be the lighter stau mass eigenstate and if one includes the latest LEP2 constraint, $m_{\tilde{\tau}} \geq 72$ GeV [12], then the solutions with squarks lighter than 600 GeV are immediately ruled out, as well as the model $n_3 = n_8 = n_H = 1$ and $n_L = n_{e^c} = n_{\nu^c} = 0$ with squarks lighter than 1.1 TeV.]

Conclusions. A notable, and somewhat remarkable, aspect of the GMSB models studied here is that although in general there are six different consistent solutions, all of them lead to similar predictions for the supersymmetric mass spectrum (see table I). The mass spectra for the supersymmetric partners (and $H^\pm$) exhibit certain characteristic features which we will now summarize: *(i)* Depending on the exact messenger content, the LSP (ignoring the possibility of light gravitino for the moment) can be either the lighter stau or the lightest neutralino. Our calculations indicate that the solution with $n_{e^c} = 1$ favors the stau as the LSP, while the one with $n_{e^c} = 0$ favors neutralino as the LSP. Our mass spectrum is dictated by the constraint of keeping $m^2_{\tilde{\tau}}$ positive and sufficiently large. *(ii)* As expected, the lighter selectron is always heavier than the lighter stau, and sleptons are always lighter than squarks, which turn out to be very heavy in this model [11], always larger than 0.6 (1.5) TeV for low (high) $\tan\beta$ solutions. [Note that the usual mass hierarchy $m_{\tilde{e}_{1,2}} \leq m_{\tilde{d}_{1,2}} \approx m_{\tilde{u}_{1,2}} \leq m_{\tilde{t}_{1,2}}$ among the masses, $m_{\tilde{l},q_{1,2}}$, of the mixed left and right sleptons and sqaurks also holds here.] *(iii)* Unlike in supersymmetric models without GMSB, here the sneutrinos always turn out to be heavier than the lighter of the charged sleptons with $m_{\tilde{\nu}_{e,\tau}} \approx m_{\tilde{\tau}_2}$. *(iv)* We find that the bilinear Higgs coupling μ lies in the $400 - 500$ GeV range (for squark masses of order 1 TeV) and can be positive *or* negative, unlike in [13]. For a vanishing B-parameter at the messenger scale μ would be positive. In general, the sign of μ does not seem to make much difference in the calculated mass spectra. However, there can be sizable effects is in, e.g., the $b \to s\gamma$ decay width, since for positive μ the interference between the SM and the chargino contributions is destructive, whereas for negative μ it is constructive. *(v)* The heavy spartner masses in these models turn out to be quite accurately directly proportional to the scale $\Lambda_{SUSY} = F/S$. [As a typical example, for the choice of $\tan\beta = 15$ and $\Lambda_{SUSY} = 50$ TeV the squark masses are around 1 TeV, while the charged Higgs boson mass is about 0.5 TeV in

all cases. The heavy sleptons, neutralinos, and charginos all have masses in the range 310-470 GeV. All of these heavy masses are linearly proportional to the scale Λ_{SUSY}.] We have computed the complete sparticle mass spectrum and the precise corresponding numerical values for all the six solutions found here will be reported in [11].

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TABLE I. Particle spectrum for a representative set of models with all possible messenger multiplicities. ($\Lambda_{\rm SUSY} = 50$TeV, $\Lambda_M = 10\Lambda_{\rm SUSY}$, $\tan\beta = 15$, sign$(\mu) = +1$)

$(n_3, n_8, n_H, n_L, n_{e^c}, n_{\nu^c})$	$\tan\beta$	μ	$m_{H^\pm}$	$m_{\tilde\chi^\pm_{1,2}}$	$m_{\tilde\chi^0_{1,2,3,4}}$	$m_{\tilde e_{1,2}}$	$m_{\tilde\tau_{1,2}}$
$\tan^2\theta_R = \alpha_{B-L}(M_R)/\alpha_R(M_R)$	$\frac{\Gamma(b\to s\gamma)}{\Gamma_{SM}}$	M_3	$m_{\tilde\nu_e}/m_{\tilde\nu_\tau}$	$m_{\tilde u_{1,2}}$	$m_{\tilde t_{1,2}}$	$m_{\tilde d_{1,2}}$	$m_{\tilde b_{1,2}}$
$(1,1,1,0,0,0)$	15	439	514	346/469	40/346/426/469	70/312	54/314
1.6	1.2	1060	302/302	1000/1045	905/1019	1001/1048	990/1006
$(1,1,0,1,0,0)$	15	441	513	347/471	40/347/429/470	90/322	78/323
1.4	1.2	1060	312/312	999/1045	904/1019	1000/1048	989/1006
$(1,1,1,0,1,0)$	15	438	514	345/468	122/346/425/468	104/318	93/320
1.6	1.2	1060	309/308	1001/1045	906/1019	1002/1048	990/1007
$(1,1,0,1,1,0)$	15	440	514	346/470	122/347/427/470	115/327	106/328
1.3	1.2	1060	317/317	1000/1045	905/1019	1001/1048	990/1007
$(1,1,1,0,0,1)$	15	438	514	345/468	40/346/425/468	100/318	89/320
1.6	1.2	1060	308/308	1001/1045	906/1019	1002/1048	990/1007
$(1,1,0,1,0,1)$	15	440	513	346/470	40/346/428/470	112/326	103/328
1.3	1.2	1060	317/316	1000/1045	905/1019	1001/1048	990/1007